

SUPERSTRINGS:THE VIEW FROM BELOW

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ABSTRACT: We review the Standard Model in a form conducive to formulating its possible short distance extensions. This depends on the value of the Higgs mass, the only unknown parameter of the model. We suggest methods to reproduce many of the small numbers in the model in terms of scale ratios, applying see-saw like ideas to the breaking of chiral symmetries. We then investigate how the $N = 1$ Standard Model extrapolated to or near the Planck scale can fit superstring models, emphasizing the use of some non-renormalizable operators generic to superstrings.

I Introduction

The Standard Model is a remarkably compact description of all fundamental matter, described in terms of only three gauge groups and nineteen parameters. Yet, it hardly looks like a fundamental theory; it probably is the low energy manifestation of a more integrated, more satisfying, and less broken-up theory. If it is indeed to be viewed as an effective low energy theory, there must exist an ultraviolet cut-off. The *raison d'être* and the value of this cut-off are the central question of fundamental theory. Since there is no experimental indication of its existence, it must be at least of the order of hundreds of GeVs. At the higher end, Nature provides us with its own cut-off, the Planck scale, the largest cut-off we can presently imagine to the standard model. However it is far removed from present experimental scales.

Local field theories of gravity certainly break down at the Planck scale, and the only known possible cure is to formulate theories which deviate from space-time locality, superstring theories. Of these, the heterotic string seems to contain all the necessary ingredients needed to reproduce the low energy world. Unfortunately, there is no detailed matching between the standard model and string theory. Yet if our world has its origin in a superstring theory, there ought to be satisfactory matching of the two at some scale below the Planck scale. The knowledge of the standard model alone, may not be sufficient to identify this match. Still, effective theories derived from superstrings typically reproduce not only renormalizable interactions, such as those found in the standard model, but also a pattern of non-renormalizable interactions, some of which may provide low energy signatures for superstring models. In addition, these theories contain not only the observed chiral families, but also a number of vector-like particles, some with electroweak quantum numbers, but with hitherto undetermined $\Delta I_W = 0$ masses. Supersymmetry at low energy provides another hint, although superstring theories do not yet predict its breaking scale.

We start by reviewing the standard model, and present arguments for low energy supersymmetry. The $N = 1$ standard model is perturbative until the Planck scale, where, we can hope to match it with superstring models. We present a mode of attack which, although

based mostly on our incomplete knowledge of the Yukawa sector, hints at the presence of certain types of non-renormalizable terms, some of which are generic to superstrings.

II The $N = 0$ Standard Model

Extension of the standard model predicts new phenomena at shorter distances, although none so far have distinguished themselves either by reproducing the *values* of the parameters, or even their multiplicity. Thus it is time to review the types of extensions which might generically explain the observed patterns.

The $N = 0$ standard model is described by three Yang-Mills groups, each with its own dimensionless gauge coupling, α_1 for the hypercharge $U(1)$, α_2 for the weak isospin $SU(2)$, and α_3 for QCD. QCD itself predicts strong CP violation, with strength proportional to a fourth dimensionless parameter $\bar{\theta}$.

The electroweak symmetry breaking Higgs sector contains two unknowns, a dimensionless Higgs self-coupling, and the Higgs mass. The “measured” value of the Fermi coupling accounts for one parameter, and the other is the value of the Higgs mass, the only parameter of the model yet to be determined from experiment. The Yukawa interactions between the fermions and the Higgs yields the nine masses of the elementary fermions. This sector also contains three mixing angles which monitor interfamily decays, and one phase which describes CP violation.

It also contains two dimensionful parameters, the Higgs mass, and the QCD confinement scale, obtained by dimensional transmutation. The QCD scale is a tiny number in Planck units $\Lambda_{QCD} \sim 10^{-20} M_{Pl}$. This small number has a natural explanation due to the logarithmic variation of the QCD coupling with scale.

The Higgs mass is unknown, but the electroweak order parameter is determined by the Fermi constant. In terms of the Planck mass it is also very small $G_F^{-1/2} \sim 10^{-17} M_{Pl}$. The origin of this small number is a matter of much speculation. In perturbation theory the Higgs mass is of the same order of magnitude as the electroweak order parameter. The most natural idea is to relate this number to dimensional transmutation associated with

new strong forces just beyond electroweak scales, called technicolor. It yields a satisfying natural explanation of this value, but these models fail to reproduce the values of the fermion masses.

Another class of theories postulates supersymmetry [1] at TeV scales. There, the electroweak order parameter is related to that of supersymmetry breaking. While not at first sight very economical, the breaking of supersymmetry automatically generates electroweak breaking [2] in a wide class of theories. The beautiful ideas of technicolor can then be applied to supersymmetry breaking, without encountering the problem of fermion masses of technicolor applied to electroweak breaking.

There are many other numbers to explain, notably in the Yukawa sector of the theory. Quark and charged lepton masses break weak isospin by half a unit, along $\Delta I_W = \frac{1}{2}$, and hypercharge by one unit, the same quantum numbers as the electroweak order parameter, which also gives the W-boson its mass. In this sense charged fermion masses should be of the same order as the W mass. This happens only for the top quark. The others are unnaturally small

$$\frac{m_{u,d}}{M_W} \sim \mathcal{O}(10^{-4}) ; \quad \frac{m_s}{M_W} \sim \mathcal{O}(10^{-3}) ; \quad \frac{m_c}{M_W} \sim \mathcal{O}(10^{-2}) ; \quad \frac{m_b}{M_W} \sim .05 .$$

Similarly for the charged leptons,

$$\frac{m_e}{M_W} \sim \mathcal{O}(10^{-5}) ; \quad \frac{m_\mu}{M_W} \sim \mathcal{O}(10^{-3}) ; \quad \frac{m_\tau}{M_W} \sim .02 ,$$

which range from the tiny to the small. Neutrino masses are predicted to be exactly zero in the standard model only because of the global chiral lepton number symmetries. However there is mounting experimental evidence that neutrinos have masses. In the absence of new degrees of freedom they are of the Majorana kind, and break weak isospin by one unit, as $\Delta I_W = 1$. Direct experimental limits on neutrino masses indicate that they are at most extremely small: $m_{\nu_e} < 10^{-17} M_W$.

The values of the three gauge parameters are known to great accuracy from measurements at low energy, although because of endemic problems associated with strong QCD,

the color coupling is the least well known. Given these parameters, we can extrapolate the standard model to shorter distances, using the renormalization group perturbatively. The most interesting effect occurs in the extrapolation of the three gauge couplings. The hypercharge and weak isospin couplings meet at a scale of 10^{13} GeV, with a value $\alpha^{-1} \approx 43$, but at that scale, the QCD coupling is much larger, $\alpha_3^{-1} \approx 38$. Thus, although the quantum numbers indicate possible unification into a larger non-Abelian group, the gauge coupling do not follow suit in this naive extrapolation. Historically, before the couplings were measured to such accuracy, it was believed that all three did indeed unify in the ultraviolet. In the ultraviolet, the values of these couplings is less disparate than at experimental scales. Similarly, nothing spectacular occurs to the Yukawa couplings. For instance, the bottom quark and τ lepton Yukawa couplings meet around 10^9 GeV, and part in the deeper ultraviolet.

The situation is potentially more extreme in the Higgs sector because of the renormalization group behavior of the Higgs self coupling [3]. We can consider two cases, depending on the value of the Higgs mass. If it is below 150 GeV, the self-coupling turns negative somewhere below Planck scale. This results in a loss of perturbation theory, with a potential unbounded from below. Using the recently announced value of the top quark mass, a Higgs mass of 120 GeV means that “instability” sets in at 1 TeV, indicating some new physics at that scale. When operative, this bound provides a low (with respect to Planck mass) energy cut-off for the standard model.

If the Higgs mass is above 200 GeV, its self-coupling rises dramatically towards its Landau pole at a relatively low energy scale. It means loss of perturbative control of the theory, and sets an upper bound on the Higgs mass since there is no evidence of any strong electroweak coupling at experimental scales. Strong coupling must happen, meaning that the Higgs is a composite. An example of this view is the technicolor scenario where the Higgs is a condensate of techniquarks.

Within a tiny range of intermediate values for the Higgs mass, the instability and triviality bounds are pushed to scales beyond the Planck length, and there is no standard

model prediction of new physics; the cut-off is indeed the Planck scale.

The dependence of the various standard model parameters on the cut-off is very different. Quantum fluctuations *additively* renormalize the Higgs mass with a term linearly proportional to the cut-off. Thus even if the Higgs mass is in a region that does not *technically* require new physics below Planck mass, its value is unnaturally small, if Planck mass is the cut-off. On the other hand, the cut-off dependence of any chiral fermion mass is only logarithmic. The reason is chiral symmetry, which is recovered by setting the fermion mass to zero. It affords a protection mechanism which results in this weak cut-off dependence.

III The $N = 1$ Standard Model

Supersymmetry avoids the *technical* naturalness problem by linking any fermion to a boson of the same mass. With exact supersymmetry, the boson mass is then protected by the chiral symmetry of the fermion. As long as supersymmetry is broken at energies in the range of TeV, this is enough protection to produce a low Higgs mass. This might seem to be small progress, since a new symmetry has been introduced to relax the strong cut-off dependence, a symmetry which has to be broken itself at a small scale, $V_{SUSY} \sim 10^{-15} M_{Pl}$.

In the $N = 1$ standard model, there are only gauge and Yukawa coupling constants. None of these couplings blow up below Planck mass. In particular, the perky Higgs self-coupling is replaced by the square of gauge and Yukawa couplings, which allows for the perturbative extrapolation all the way to Planck scale, opening the way for comparison with fundamental theory!

There are tantalizing hints of simplicity in the extrapolation of the couplings. Firstly the gauge couplings seem to be much closer to unification, and at a scale large enough not to be invalidated by proton decay bounds. The hypercharge and weak isospin couplings meet at a scale of the order of 10^{16} GeV, with a value $\alpha^{-1} \approx 25$, and the QCD coupling is much closer to, if not right on the same value [4]. It may still be a shade higher than the others, with $(\alpha^{-1} - \alpha_3^{-1}) \leq 1.5$.

The second remarkable thing is that with simple boundary conditions at or near Planck mass, inspired by a simple picture of supersymmetry breaking, the renormalization group drives one of the Higgs masses to imaginary values in the infrared. This in turns triggers electroweak breaking[2], made possible by the large top quark mass!

The Higgs mass is not arbitrarily high in the minimal extension. At tree-level, it is predicted to be below the Z-mass, but it suffers large radiative corrections due to the top Yukawa coupling, raising it above the Z, but not by an arbitrarily large amount [5].

This general scheme allows us to study the pattern of fermion masses at these shorter distances; there are more regularities with supersymmetry. For instance, the bottom quark and τ masses seem to unify at or around 10^{16-17} GeV [6], the same scale where the gauge couplings converge.

The most striking aspect of the fermion masses is that only one chiral family has large masses, leading us to consider theories where the tree-level Yukawa matrices are simply of the form

$$\mathbf{Y}_{u,d,e} = y_{t,b,\tau} \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix} .$$

These matrices imply a global chiral symmetry, $U(2)_L \times U(2)_R$, in each charged sector. The hierarchy between the bottom and top quark masses requires explanation. In the $N = 1$ model, it is linked to another parameter which comes from the Higgs sector, the ratio of the *vev* of the two Higgs. Hence it may not pertain to properties of the Yukawa matrices. Why are the other two families so light? Starting from the rank two Yukawa matrices, we must find a scheme by which the zeros get filled, presumably in higher orders of perturbation theory. In order to see how this might come about, let us examine one well-known case in which small numbers are naturally generated, the see-saw mechanism [7].

In the standard model, the neutrino Majorana mass matrix is zero at tree-level. A detailed examination shows that these zeros are protected from quantum corrections by conservation of chiral global lepton number for each species. In the see-saw mechanism,

the usual neutrinos are mixed with new electroweak singlet fields (neutral leptons), by $\Delta I_w = 1/2$ terms, of electroweak breaking strength, to give them the same lepton numbers. These new particles can acquire $\Delta I_W = 0$ Majorana masses, M , of any magnitude, in particular well above the electroweak scale. Upon diagonalization, this generates a mass for the familiar neutrinos, depressed from typical electroweak values by the ratio of scale $\frac{m}{M}$, where m is the electroweak order parameter. A scale ratio between electroweak and chiral lepton number breaking is used to generate a small number.

A similar analysis may apply to the charged Yukawa matrices, where the zeros are also protected by chiral symmetries. We first couple the massless fermions of the first two families to new fermions with similar quantum numbers, thereby extending the chiral symmetries to them. Unlike the neutral case, these new fermions have electroweak charges, and cannot have Majorana masses. Breaking of chiral symmetry is done by their Dirac masses, which requires the presence of vector-like partners (this differs from the neutral sector), along the $\Delta I_W = 0$ direction at a new undetermined scale M .

One may also take the point of view that these non-renormalizable operators come from physics beyond the Planck scale, in which case, the question is relegated to one of classifying the possible non-renormalizable operators.

Now that low energy supersymmetry allows its perturbative extrapolation deep in the ultraviolet, we may ask how it can be made to match with more fundamental theories, one type being superstrings.

IV. Matching to Superstrings

Superstring theories are not understood in detail, but some of their generic features are evident. We are interested in the effective theory they generate at or near Planck scale. An estimate of string effects indicates that this scale is related to the gauge coupling through the formula [[8]]

$$M_U \approx 2.5 \sqrt{\alpha_U} \times 10^{18} \text{ GeV} .$$

With $M_X = 10^{16} \text{ GeV}$, and $\alpha_U^{-1} < \alpha_X^{-1} \approx 25$, this implies that contact with the superstring

can be made provided that $M_U/M_X > 50$: there is a discrepancy in the matching of scales.

The second feature of superstring theories is to produce at lower energies remnants of **27** and $\overline{\mathbf{27}}$ representations of E_6 . In the effective low energy theory, these yield the three chiral families, together with many vector-like particles, capable of sharing quantum numbers with the chiral particles. These particles may be used in see-saw like mechanisms to generate small numbers in the Yukawa matrices. It also means that there may be many intermediate thresholds between the supersymmetry and unification scales

A third feature of superstring theories is that the gauged group at the string scale is larger than the standard model group. This implies the existence of more gauge bosons at intermediate scales and many vector-like particles with electroweak quantum numbers. The values of their masses to be determined by the flat directions in the superpotential, and the discrete symmetries of the particular model.

A fourth generic feature is the existence of a local $U(1)$ symmetry, with anomaly cancelled through the Green-Schwarz mechanism. This symmetry is however broken close to the Planck scale. Does any trace of this symmetry appear in the extrapolated low energy standard model? Ibàñez [9] has argued that this symmetry can be used to fix the Weinberg angle in superstring theories. Following Ibàñez and Ross[10], we argue[11] that this Abelian symmetry sets the dimensions of the Froggatt and Nielsen[12] Yukawa operators. Are any of these features present in the extrapolated low energy theory?

Consider first the unification of the gauge couplings. It is predicated on two assumptions: that the weak hypercharge coupling is normalized to its unification into a higher rank Lie group, such as $SU(5)$, $SO(10)$ or E_6 [13]), and on the absence of intermediate thresholds with matter carrying strong or electroweak quantum numbers between 1 TeV and 10^{16} GeV. The gauge couplings may not exactly unify at M_X , and we may want to alter this simple picture by requiring at least one intermediate threshold between the SUSY scale and the illusory unification scale at M_X to obtain unification at the string scale M_U [14]].

At one-loop, the couplings $\alpha_i^{-1}(t)$ for the three gauge groups, ($i = 1, 2, 3$ for $U(1)_Y$,

$SU(2)_L$, and $SU(3)^c$, respectively) run with scale according to

$$= \alpha_i^{-1}(t_X) + \frac{b_i}{2\pi}(t - t_X) ,$$

where

$$t = \ln(\mu/\mu_0) , \quad t_X = \ln(M_X/\mu_0) ,$$

and μ_0 is an arbitrary reference energy. For the three families and two Higgs doublets of the minimal supersymmetric standard model, we have

$$b_1 = -\frac{33}{5} ; \quad b_2 = -1 ; \quad b_3 = 3 .$$

Since the low energies values of α_1 and α_2 are known with the greatest accuracy, we use their trajectories to define t_X as the scale at which they meet:

$$\alpha_X^{-1} \equiv \alpha_1^{-1}(t_X) = \alpha_2^{-1}(t_X) .$$

The extrapolated data show that $\alpha_X^{-1} \approx 25$, with $M_X \approx 10^{16}$ GeV. We do not assume precisely the same value for $\alpha_3(t_X)$ at that scale; rather we set

$$\alpha_X^{-1} = \alpha_3^{-1}(t_X) + \Delta .$$

Present uncertainties on the QCD coupling suggest that

$$|\Delta| \leq 1.5 .$$

Suppose there is an intermediate threshold above supersymmetry at

$$t_I = \ln(M_I/\mu_0) ; \quad t_I < t_X ,$$

caused by new vector-like particles with electroweak singlet masses at M_I . Their effect is to alter the b_i coefficients:

$$b_i \rightarrow b_i - \delta_i , \quad i = 1, 2, 3 .$$

By requiring unification at M_U , we find the constraints

$$\frac{r}{14} = \frac{t_U - t_X}{t_U - t_I}$$

and

$$\frac{q}{4} = \frac{t_U - t_X - \pi\Delta/2}{t_U - t_I} ,$$

written in terms of

$$q \equiv \delta_3 - \delta_2 \quad \text{and} \quad \frac{2}{5}r \equiv \delta_2 - \delta_1 .$$

For vector-like matter generated from superstrings, q and r are integers. The value of the gauge coupling at unification is now

$$\alpha_U^{-1} = \alpha_X^{-1} - \frac{1}{2\pi} [\delta_2(t_U - t_I) + t_U - t_X] .$$

These equations have solutions for non-exotic matter. For instance when $\Delta = 0.82$ with $r = 5$, $q = 1$, we get

$$M_U = 7.5 \times 10^{17} \text{GeV} ; \quad M_I = 4.4 \times 10^{12} \text{GeV} ; \quad \alpha_U^{-1} = 11 .$$

However most solutions do not allow large M_X/M_I .

In realistic superstring models, the assumption of one intermediate scale is probably not justified. For several intermediate thresholds, by applying these equations repeatedly, we obtain similar equations, with q and r replaced by average quantities which are no longer integers. Take for instance the interesting example of the 3-family Calabi-Yau superstring model of ref. [15]. After flux breaking, the surviving gauge group is

$$SU(3)_L \times SU(3)^c \times SU(3)_R .$$

There are at least two *a priori* distinct intermediate scale order parameters associated with each reduction in rank. Many chiral superfields survive flux breaking: 9 leptons, 6 mirror leptons, 7 quarks, 4 mirror quarks, 7 antiquarks, 4 mirror antiquarks. With all these particles concentrated at one mass scale, there is no solution, but this is hardly realistic. It is convenient to define the effective intermediate scale as the average intermediate scale weighted by δ_2 , i.e.,

$$t_{\bar{I}} \equiv \frac{\sum_{a=1}^N t_{Ia} \delta_{2a}}{\sum_{a=1}^N \delta_{2a}} = \frac{1}{29} \sum_{a=1}^N t_{Ia} \delta_{2a} .$$

Taking 5×10^{17} GeV as a minimum for M_U , we find the high value.

$$M_{\overline{T}} > 3 \times 10^{15} \text{ GeV}$$

Gauge coupling unification can be attained in this example in a calculably perturbative way, but it requires that many of the vector-like particle which survive flux breaking be very close to the string scale, and that electroweak-doublet vectorlike particles be heavier on average than the strongly-interacting electroweak singlet vectorlike particles.

It might seem rather surprising that in the MSSM the gauge couplings should appear to be nicely headed for unification at M_X , only to be redirected to a new meeting place at M_U , but the apparent perverseness of this situation allows us put some non-trivial constraints on the scenario.

Let us now turn to the last topic, the possibility of an Abelian gauge symmetry, with anomaly cancelled by the Green-Schwarz mechanism. Such can be recognized if it plays a role in determining the dimensions of the entries of the Yukawa matrices[10,11]. The most general Abelian charge that can be assigned to the particles of the Minimal Supersymmetric Standard Model, with μ term, can be written as

$$X = X_0 + X_3 + \sqrt{3}X_8 ,$$

where X_0 is the family independent part, X_3 is along λ_3 , and X_8 is along λ_8 . We set

$$X_{3,8} = (a_{3,8}, b_{3,8}, c_{3,8}, d_{3,8}, e_{3,8}) ,$$

where the entries correspond to the components in the family space of the fields \mathbf{Q} , $\overline{\mathbf{u}}$, $\overline{\mathbf{d}}$, L , and \overline{e} , respectively. Both Higgs doublets have the same zero X-charge, without loss of generality, since an imbalance can be created by mixing in the hypercharge Y .

With the tree-level Yukawa coupling *only* to the third family, we obtain the constraints

$$\frac{m+n}{3} = 2(a_8 + b_8) , \quad \frac{m+p}{3} = 2(a_8 + c_8) , \quad \frac{q+r}{3} = 2(d_8 + e_8) .$$

The excess X-charge at each of their entries is made up by powers of an electroweak singlet field, resulting in operators of higher dimensions. A typical term would be of the form

$$\mathbf{Q}_i \overline{\mathbf{u}}_j H_u \left(\frac{\theta}{M} \right)^{n_{ij}} ,$$

where θ is some field with charge x , M is some large scale, and n_{ij} is the excess X-charge listed above for the Yukawa matrices. The exponents are determined by X-symmetry. In order to produce a small coefficient, the i th and j th fermions need to go through a number of intermediate steps to interact. The larger the number steps, the larger n_{ij} , and the smaller the entry in the effective Yukawa matrix. This approach was advocated long ago by Froggatt and Nielsen [12]. This yields approximate zeros in the matrices, creating textures[16] . For example, in the charge 2/3 sector,

$$n_{12}x = 3(a_8 + b_8) + a_3 - b_3 .$$

Since θ may have a large expectation value, it is likely accompanied by its vector-like partner $\bar{\theta}$, with opposite charge, showing that the exponents n_{ij} need not be positive, but if all the n_{ij} are positive, several interesting phenomenological consequences follow[[11]]. First the n_{ij} exponents are not all independent, resulting in order of magnitude estimates among the Yukawa matrix elements

$$(Y)_{11} \sim \frac{(Y)_{13}(Y)_{31}}{(Y)_{33}} ,$$

$$(Y)_{22} \sim \frac{(Y)_{23}(Y)_{32}}{(Y)_{33}} ,$$

valid for each of the three charge sectors. These relations are consistent with many of the allowed textures. Another important consequence is that the X-charge of the determinant in each charge sector is *independent* of the texture coefficients that distinguish between the two lightest families

$$\text{charge } \frac{2}{3} : 6(a_8 + b_8) \equiv U , \text{ charge } -\frac{1}{3} : 6(a_8 + c_8) \equiv D , \text{ charge } -1 : 6(d_8 + e_8) \equiv L .$$

Let the value of $\frac{\theta}{M}$ be a small parameter λ . In the simplest case, this parameter would be the same for all three charge sectors. Then we have

$$\frac{m_d m_s m_b}{m_e m_\mu m_\tau} \sim \mathcal{O}(\lambda^{(D-L)/x}) .$$

It is more difficult to compare the up and down sectors in this way since we do not know the value of $\tan \beta$, which sets the normalization between the two sectors

$$\frac{m_u m_c m_t}{m_d m_s m_b} \sim \tan^3 \beta \times \mathcal{O}(\lambda^{(U-D)/x}) .$$

Since this ratio is much larger than one, it means either that $\tan\beta$ is itself large, with U close to D , or that $\tan\beta$ is not large, but $D > U$.

In general, the X symmetry is anomalous. The three chiral families contribute to the mixed gauge anomalies as follows

$$\begin{aligned} C_3 &= 2m + n + p , \\ C_2 &= 3m + q + 2s , \\ C_1 &= \frac{1}{3}m + \frac{8}{3}n + \frac{2}{3}p + q + 2r + 2s . \end{aligned}$$

The subscript denotes the gauge group of the Standard Model, *i.e.* $1 \sim U(1)$, $2 \sim SU(2)$, and $3 \sim SU(3)$. The X-charge also has a mixed gravitational anomaly, which is simply the trace of the X-charge,

$$C_g = (6m + 3n + 3p + 2q + r + 4s) + C'_g ,$$

where C'_g is the contribution from the particles that do not appear in the minimal $N = 1$ model. The last anomaly coefficient is that of the X-charge itself, C_X , which is the sum of the cubes of the X-charge.

It was suggested by Ibàñez[9], that an anomalous $U(1)$ symmetry, with its anomalies cancelled through the Green-Schwarz mechanism, is capable of relating the ratio of gauge couplings to the ratios of anomaly coefficients

$$\frac{C_i}{k_i} = \frac{C_X}{k_X} = \frac{C_g}{k_g} ,$$

which relates the Weinberg angle to the anomaly coefficients, without the use of Grand Unification. The k_i are the Kac-Moody levels; are integers for the non-Abelian factors only. The mixed YXX anomaly, however, must vanish by itself.

We demand that the non-Abelian gauge groups have the same Kac-Moody levels, which means that

$$C_2 = C_3 \quad \text{or} \quad q = n + p - m - 2s .$$

Secondly we require that at or near the unification or string scale, the Weinberg angle have the value

$$\sin^2 \theta_W = \frac{3}{8} ,$$

which translates into the further constraint

$$5C_2 = 3C_1 \quad \text{or} \quad r = 2m - n .$$

These equations are sufficient to infer that $L = D$, which implies, remarkably enough, that the products of the charged lepton masses is of the same order of magnitude as that of the down-type quarks[11]. It is satisfying to note that extrapolation of the masses to short distances indicates that these two products are in fact roughly equal in the deep ultraviolet.

This formalism has been used[10] to generate symmetric textures, of the kind found to be allowed by experiment[16]. Work is in progress to determine how these equations constrain possible textures. One result is that it appears to be difficult to generate acceptable constraints, without invoking Green-Schwarz cancellation. In that case, this particular way of generating textures would require the type of mechanism that is generic to superstrings!

The following examples have shown how several problems with the standard model might require a superstring explanation. While it is clearly too soon to claim to have made the connection, it is a fruitful path to take, as we must try to match the apparent unruliness and ugliness of the world we observe to the more beautiful and satisfying constructs of our imagination. Feza would have liked that.

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References

- 1) For reviews, see H. P. Nilles, *Phys. Rep.* **110** (1984) 1 and H. E. Haber and G. L. Kane, *Phys. Rep.* **117** (1985) 75.
- 2) L. E. Ibáñez and G. G. Ross, *Phys. Lett.* **110B** (1982) 215; K. Inoue, A. Kakuto, H. Komatsu, and S. Takeshita, *Prog. Theor. Phys.* **68** (1982) 927; L. Alvarez-Gaumé, M. Claudson, and M. Wise, *Nucl. Phys.* **B207** (1982) 16; J. Ellis, J. S. Hagelin, D. V. Nanopoulos, and K. Tamvakis, *Phys. Lett.* **125B** (1983) 275.

- 3) M. Sher, *Phys. Rep.* **179** (1989) 273.
- 4) U. Amaldi, W. de Boer, and H. Furstenau, *Phys. Lett.* **B260** (1991) 447; J. Ellis, S. Kelley and D. Nanopoulos, *Phys. Lett.* **260B** (1991) 131; P. Langacker and M. Luo, *Phys. Rev.* **D44** (1991) 817.
- 5) G. L. Kane, C. Kolda, and J. D. Wells, *Phys. Rev. Lett.* **70** (1993) 2686.
- 6) H. Arason, D. J. Castaño, B. Keszthelyi, S. Mikaelian, E. J. Piard, P. Ramond, and B. D. Wright, *Phys. Rev. Lett.* **67** (1991) 2933; A. Givon, L. J. Hall, and U. Sarid, *Phys. Lett.* **271B** (1991) 138.
- 7) M. Gell-Mann, P. Ramond, and R. Slansky in Sanibel Talk, CALT-68-709, Feb 1979, and in *Supergravity* (North Holland, Amsterdam 1979). T. Yanagida, in *Proceedings of the Workshop on Unified Theory and Baryon Number of the Universe*, KEK, Japan, 1979.
- 8) J.P. Derendinger, S. Ferrara, C. Kounnas, and F. Zwirner, *Nucl. Phys.* **B372** (1992) 145.
- 9) L. Ibáñez, *Phys. Lett.* **B303** (1993) 55.
- 10) L. Ibáñez and G. G. Ross, *Phys. Lett.* **B332** (1994) 100.
- 11) P. Binétruy and P. Ramond, in preparation.
- 12) C. Froggatt and H. B. Nielsen *Nucl. Phys.* **B147** (1979) 277.
- 13) J. C. Pati and A. Salam, *Phys. Rev.* **D10** (1974) 275; H. Georgi and S. Glashow, *Phys. Rev. Lett.* **32** (1974) 438; H. Georgi, in *Particles and Fields-1974*, edited by C.E. Carlson, AIP Conference Proceedings No. 23 (American Institute of Physics, New York, 1975) p.575; H. Fritzsch and P. Minkowski, *Ann. Phys. NY* **93** (1975) 193; F. Gürsey, P. Ramond, and P. Sikivie, *Phys. Lett.* **60B** (1975) 177.
- 14) S. Martin and P. Ramond, in preparation.
- 15) B. R. Greene, K. H. Kirklin, P. J. Miron, and G. G. Ross, *Phys. Lett.* **B180** (1986) 69; *Nucl. Phys.* **B278** (1986) 667; *Nucl. Phys.* **B292** (1987) 606.

- 16) P. Ramond, R.G. Roberts and G. G. Ross, *Nucl. Phys.* **B406** (1993) 19.